

Leading order covariant hyperon-nucleon interactions in chiral effective field theory

Kai-Wen Li^a, Xiu-Lei Ren^b, Li-Sheng Geng^{a,c,*}, Bingwei Long^d

^a*School of Physics and Nuclear Energy Engineering and International Research Center for Nuclei and Particles in the Cosmos, Beihang University, Beijing 100191, China*

^b*School of Physics and State Key Laboratory of Nuclear Physics and Technology, Peking University, Beijing 100871, China*

^c*Beijing Key Laboratory of Advanced Nuclear Materials and Physics, Beihang University, Beijing 100191, China*

^d*Center for Theoretical Physics, Department of Physics, Sichuan University, 29 Wang-Jiang Road, Chengdu, Sichuan 610064, China*

Abstract

We apply a recently proposed covariant power counting to study strangeness $S = -1$ $\Lambda N - \Sigma N$ interactions in chiral effective field theory. At leading order, Lorentz invariance introduces 12 low energy constants in contrast to the Weinberg counting based on naïve dimensional analysis, where only five appear. The Kadyshevsky equation is adopted to resum the potential in order to account for the non-perturbative nature of hyperon-nucleon interactions. A fit to the 36 hyperon-nucleon scattering data yields $\chi^2 \simeq 17$, which is comparable with the sophisticated phenomenological models and the next-to-leading order Weinberg approach. The results hint that a more efficient formulation of the hyperon-nucleon interaction can be achieved in covariant chiral effective field theory, in agreement with the results in the nucleon-nucleon sector using a similar framework.

Keywords: Hyperon-nucleon interactions, covariant chiral effective field theory, Kadyshevsky equation

1. Introduction

Since the quantum number *strangeness* was introduced [1, 2] and the first observation of Λ hypernuclei [3] in 1953, strangeness nuclear physics has always been at the frontier of experimental and theoretical nuclear physics. In recent years, open questions such as the charge symmetry breaking in $A = 4$ Λ -hypernuclei [4] and the existence of H-dibaryon [5] have attracted a lot of attention [6–14]. In facilities like J-PARC in Japan and FAIR in Germany, many important studies are being pursued, e.g., the level spectra and decay properties of Λ , double Λ and Ξ hypernuclei [8, 15], the $\Sigma^+ p$ scattering [16], and the final state interactions in production reactions, such as $\bar{p}p \rightarrow K^+ \Lambda p$ [17], which can provide information on the ΛN scattering lengths. Meanwhile, theoretical few- and many-body calculations of hypernuclei have made steady progress, see, e.g., Refs. [18, 19]. One particularly interesting ongoing issue is about the role of hyperons in the core of neutron stars, known as the *hyperon puzzle*: Nuclear many-body calculations incorporating hyperon degrees of freedom [20–24] have difficulties obtaining a two-solar mass neutron star that was recently observed [25, 26].

As the most important theoretical input for few- and many-body calculations, baryon-baryon interactions play an indispensable role in studies of hypernuclear physics. Although many efforts have been made to derive them, previous theoretical investigations were mainly based on phenomenological meson-exchange models [27–34] and quark models [35–41]. In the past two decades, two breakthroughs have occurred

in constructing model-independent baryon-baryon interactions, and both of them are more closely related to Quantum Chromodynamics (QCD), the underlying theory of strong interactions. One is lattice QCD simulations [42–47], which provide an *ab initio* numerical solution to QCD from first principles. With the ever-growing of computing power and evolving numerical algorithms, lattice QCD simulations are approaching the physical world [48, 49], thus providing us with more information and constraints on baryon-baryon interactions. The other is chiral effective field theory (χ EFT), which has achieved great successes in nucleon-nucleon (NN) interactions [50–52] following Weinberg’s proposal [53, 54]. The latter approach has been generalized to antinucleon-nucleon [55], hyperon-nucleon (YN) [56–58] and multi-strangeness systems [59–61]. The main advantage of χ EFT is that by using a power counting scheme, one can improve calculations systematically by going to higher orders in powers of external momenta and light quark masses, and estimating the uncertainties of any given order. Furthermore, three- and four-body forces automatically arise as we push through the hierarchy of chiral forces.

However, the Weinberg approach for baryon-baryon interaction is based on a non-relativistic formalism. It suffers from high sensitivities to ultraviolet cutoffs, that is, renormalization group invariance is violated, risking severe model dependence of short-range physics. Various opinions on this issue can be found in Refs. [62–69]. In two recent papers, Epelbaum and Gegelia have proposed a new approach (referred to as the EG approach in the present letter) to NN scattering in χ EFT [70, 71], where the time-ordered perturbation theory is employed and relativistic effects are partially retained. At leading order (LO), the potential remains unchanged but the scattering equation changes to the Kadyshevsky equation, compared to

*Corresponding Author

Email address: lisheng.geng@buaa.edu.cn (Li-Sheng Geng)

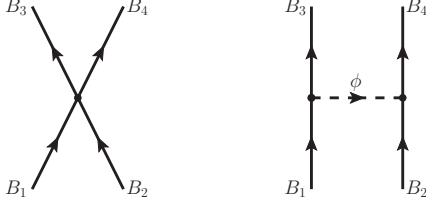


Figure 1: Nonderivative four-baryon contact terms and one-pseudoscalar-meson exchanges at LO. The solid lines denote incoming and outgoing baryons ($B_{1,2,3,4}$), and the dashed line denotes the exchanged pseudoscalar meson ϕ .

the Lippmann-Schwinger equation with nonrelativistic nucleon propagators in the Weinberg approach. Although this turned out to improve the description of the Nijmegen partial wave analysis (PWA) [72], a higher order contact term was still needed in 3P_0 partial wave to achieve renormalization group invariance. We applied the EG approach to the strangeness $S = -1$ YN system [73] and found that the best description of the experimental data is quantitatively similar to that of the Weinberg approach, and that cutoff dependence is mitigated but not removed.

Partly motivated by the successes of covariant χ EFT in the one-baryon and heavy-light systems [74–82], a new covariant power counting is proposed in Ref. [83] to study NN chiral interactions. The covariant treatment of baryons maintains all the symmetries and analyticities, and, at LO, it results in a more efficient scheme in the sense that it requires fewer low energy constants (LECs) to achieve similar accuracy, in comparison with the next-to-leading order (NLO) Weinberg scheme. In the present letter, we apply this scheme to YN scattering with strangeness $S = -1$. Given the fact that experimental data are scarce in the YN sectors, a more efficient formulation of baryon-baryon forces with fewer LECs is particularly relevant.

2. Formalism

In the covariant power counting [83], the full baryon spinor is retained to maintain Lorentz invariance

$$u_B(\mathbf{p}, s) = N_p \left(\frac{1}{\epsilon_B} \right) \chi_s, \quad N_p = \sqrt{\frac{\epsilon_B}{2M_B}}, \quad (1)$$

where $\epsilon_B = E_B + M_B$ and $E_B = \sqrt{\mathbf{p}^2 + M_B^2}$, while a non-relativistic reduction of u_B is employed in the Weinberg approach.

According to naive dimensional analysis, the LO baryon-baryon interactions include nonderivative four-baryon contact terms (CT) and one-pseudoscalar-meson exchange (OME) potentials, as shown in Fig. 1,

$$V_{\text{LO}} = V_{\text{CT}} + V_{\text{OME}}. \quad (2)$$

2.1. Four-baryon contact terms

The Lagrangian terms for nonderivative four-baryon contact interactions are

$$\begin{aligned} \mathcal{L}_{\text{CT}}^1 &= \frac{C_1^1}{2} \text{tr} \left(\bar{B}_a \bar{B}_b (\Gamma_i B)_b (\Gamma_i B)_a \right), \\ \mathcal{L}_{\text{CT}}^2 &= \frac{C_2^2}{2} \text{tr} \left(\bar{B}_a (\Gamma_i B)_a \bar{B}_b (\Gamma_i B)_b \right), \\ \mathcal{L}_{\text{CT}}^3 &= \frac{C_3^3}{2} \text{tr} \left(\bar{B}_a (\Gamma_i B)_a \right) \text{tr} \left(\bar{B}_b (\Gamma_i B)_b \right), \end{aligned} \quad (3)$$

where tr indicates trace in flavor space (u , d , and s); Γ_i are the elements of the Clifford algebra,

$$\Gamma_1 = 1, \quad \Gamma_2 = \gamma^\mu, \quad \Gamma_3 = \sigma^{\mu\nu}, \quad \Gamma_4 = \gamma^\mu \gamma_5, \quad \Gamma_5 = \gamma_5; \quad (4)$$

and C_i^m ($m = 1, 2, 3$) are the LECs corresponding to independent four-baryon operators. The ground-state octet baryons are collected in the 3×3 traceless matrix:

$$B = \begin{pmatrix} \frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda}{\sqrt{6}} & \Sigma^+ & p \\ \Sigma^- & -\frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda}{\sqrt{6}} & n \\ \Xi^- & \Xi^0 & -\frac{2\Lambda}{\sqrt{6}} \end{pmatrix}. \quad (5)$$

The potentials corresponding to the contact interaction can be written as

$$V_{\text{CT}}^{BB'} = C_i (\bar{u}_3 \Gamma_i u_1) (\bar{u}_4 \Gamma_i u_2). \quad (6)$$

They are first calculated in helicity basis and then transformed to $|LSJ\rangle$ basis [32]. We found that they contribute to all partial waves that have total angular momentum $J \leq 1$. We choose the LECs in 1S_0 , 3S_1 and 3P_1 to be independent, as shown below. The partial wave projected potentials are

$$\begin{aligned} V_{\text{CT}}^{BB'}(^1S_0) &= \xi_B \left[(C_1 + C_2 - 6C_3 + 3C_4)(1 + R_p^2 R_{p'}^2) \right. \\ &\quad \left. + (3C_2 + 6C_3 + C_4 + C_5)(R_p^2 + R_{p'}^2) \right] \\ &\equiv \xi_B \left[C_{1S_0}^{BB'}(1 + R_p^2 R_{p'}^2) + \hat{C}_{1S_0}^{BB'}(R_p^2 + R_{p'}^2) \right], \end{aligned} \quad (7)$$

$$\begin{aligned} V_{\text{CT}}^{BB'}(^3S_1) &= \xi_B \left[\frac{1}{9}(C_1 + C_2 + 2C_3 - C_4)(9 + R_p^2 R_{p'}^2) \right. \\ &\quad \left. + \frac{1}{3}(C_2 + 2C_3 - C_4 - C_5)(R_p^2 + R_{p'}^2) \right] \\ &\equiv \xi_B \left[\frac{1}{9}C_{3S_1}^{BB'}(9 + R_p^2 R_{p'}^2) + \frac{1}{3}\hat{C}_{3S_1}^{BB'}(R_p^2 + R_{p'}^2) \right], \end{aligned} \quad (8)$$

$$\begin{aligned} V_{\text{CT}}^{BB'}(^3P_1) &= \xi_B \left[-\frac{4}{3}(C_1 - 2C_2 + 4C_3 + 2C_4 - C_5)R_p R_{p'} \right] \\ &\equiv \xi_B \left[-\frac{4}{3}C_{3P_1}^{BB'}R_p R_{p'} \right], \end{aligned} \quad (9)$$

$$\begin{aligned} V_{\text{CT}}^{BB'}(^3P_0) &= \xi_B \left[-2(C_1 - 4C_2 - 4C_4 + C_5)R_p R_{p'} \right] \\ &= \xi_B \left[-2(-C_{1S_0}^{BB'} - \hat{C}_{1S_0}^{BB'} + 2C_{3S_1}^{BB'} - 2\hat{C}_{3S_1}^{BB'})R_p R_{p'} \right], \end{aligned} \quad (10)$$

$$V_{\text{CT}}^{BB'}(^1P_1) = \xi_B \left[-\frac{2}{3}(C_1 + C_5)R_p R_{p'} \right] \\ = \xi_B \left[-\frac{2}{3}(C_{3S1}^{BB'} - \hat{C}_{3S1}^{BB'})R_p R_{p'} \right], \quad (11)$$

$$V_{\text{CT}}^{BB'}(^3S_1 - ^3D_1) = \xi_B \left[\frac{2}{9} \sqrt{2}(C_1 + C_2 + 2C_3 - C_4)R_p^2 R_{p'}^2 \right. \\ \left. + \frac{2}{3} \sqrt{2}(C_2 + 2C_3 - C_4 - C_5)R_p^2 \right] \\ = \xi_B \left[\frac{2}{9} \sqrt{2}C_{3S1}^{BB'}R_p^2 R_{p'}^2 + \frac{2}{3} \sqrt{2}\hat{C}_{3S1}^{BB'}R_p^2 \right], \quad (12)$$

$$V_{\text{CT}}^{BB'}(^3D_1 - ^3S_1) = \xi_B \left[\frac{2}{9} \sqrt{2}(C_1 + C_2 + 2C_3 - C_4)R_p^2 R_{p'}^2 \right. \\ \left. + \frac{2}{3} \sqrt{2}(C_2 + 2C_3 - C_4 - C_5)R_p^2 \right] \\ = \xi_B \left[\frac{2}{9} \sqrt{2}C_{3S1}^{BB'}R_p^2 R_{p'}^2 + \frac{2}{3} \sqrt{2}\hat{C}_{3S1}^{BB'}R_p^2 \right], \quad (13)$$

$$V_{\text{CT}}^{BB'}(^3D_1) = \xi_B \left[\frac{8}{9}(C_1 + C_2 + 2C_3 - C_4)R_p^2 R_{p'}^2 \right] \\ = \xi_B \left[\frac{8}{9}C_{3S1}^{BB'}R_p^2 R_{p'}^2 \right], \quad (14)$$

where $\xi_B = N_p^2 N_{p'}^2$, $R_p = |\mathbf{p}|/\epsilon_p$, $R_{p'} = |\mathbf{p}'|/\epsilon_{p'}$ and $M_B = 1080$ MeV stands for the SU(3) average mass of the octet baryons. \mathbf{p} and \mathbf{p}' denote the initial and final momenta, respectively. To recover the potentials in the Weinberg approach we simply take $R_p = R_{p'} = 0$ and $\xi_B = 1$. Assuming strict SU(3) symmetry, one can easily obtain the contact terms for the $S = -1$ hyperon-nucleon system. Finally we have three Feymann diagrams for the contact terms, distinguished by different baryon species, and 12 independent LECs: $C_{150}^{\Lambda\Lambda}$, $\hat{C}_{150}^{\Lambda\Lambda}$, $C_{150}^{\Sigma\Sigma}$, $\hat{C}_{150}^{\Sigma\Sigma}$, $C_{3S1}^{\Lambda\Lambda}$, $\hat{C}_{3S1}^{\Lambda\Lambda}$, $C_{3S1}^{\Sigma\Sigma}$, $\hat{C}_{3S1}^{\Sigma\Sigma}$, $C_{3S1}^{\Lambda\Sigma}$, $\hat{C}_{3S1}^{\Lambda\Sigma}$, $C_{3P1}^{\Lambda\Lambda}$, $\hat{C}_{3P1}^{\Lambda\Lambda}$. The superscript $Y_1 Y_2$ denotes the hyperons in the reaction of $Y_1 N \rightarrow Y_2 N$, as used in Ref. [56].

2.2. OME potentials

The OME potentials are derived from the covariant SU(3) meson-baryon Lagrangian,

$$\mathcal{L}_{MB}^{(1)} = \text{tr} \left(\bar{B}(i\gamma_\mu D^\mu - M_B)B \right. \\ \left. - \frac{D}{2} \bar{B} \gamma^\mu \gamma_5 \{u_\mu, B\} - \frac{F}{2} \bar{B} \gamma_\mu \gamma_5 [u_\mu, B] \right), \quad (15)$$

where $D^\mu B = \partial_\mu B + [\Gamma_\mu, B]$ and D and F are the axial vector couplings. In the numerical analysis, we use $D + F = g_A = 1.26$ and $F/(F + D) = 0.4$, where g_A is the nucleon axial vector coupling constant. Γ_μ and u_μ are the vector and axial vector combinations of the pseudoscalar-meson fields and their derivatives,

$$\Gamma_\mu = \frac{1}{2} [u^\dagger \partial_\mu u + u \partial_\mu u^\dagger], \quad u_\mu = i(u^\dagger \partial_\mu u - u \partial_\mu u^\dagger),$$

where $u^2 = U = \exp(i\frac{\sqrt{2}\phi}{f_0})$, with the pseudoscalar-meson decay constant $f_0 = 93$ MeV [58], and the traceless matrix ϕ collecting the pseudoscalar-meson fields:

$$\phi = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & K^0 \\ K^- & \bar{K}^0 & -\frac{2\eta}{\sqrt{6}} \end{pmatrix}. \quad (16)$$

The potentials for OME can be expressed in a generic form:

$$V_{\text{OME}} = -N_{B_1 B_3 \phi} N_{B_2 B_4 \phi} \frac{(\bar{u}_3 \gamma^\mu \gamma_5 q_\mu u_1)(\bar{u}_4 \gamma^\nu \gamma_5 q_\nu u_2)}{q^2 - m^2} \mathcal{I}_{B_1 B_2 \rightarrow B_3 B_4}, \quad (17)$$

where $q = p' - p$ is the momentum transfer, the SU(3) coefficient $N_{BB'\phi}$ and isospin factor $\mathcal{I}_{B_1 B_2 \rightarrow B_3 B_4}$ are listed in Refs. [56, 73], as well as seven corresponding Feynman diagrams for strangeness $S = -1$. One can easily obtain V_{OME} in the $|LSJ\rangle$ basis following the same procedure as that for the contact terms. We note that by the mass differences of the exchanged mesons the SU(3) symmetry is broken.

2.3. Scattering equation

The infrared enhancement in two-baryon propagations gives the theoretical argument for low-energy baryon-baryon interactions to be nonperturbative [54]. As a result, one needs to iterate the potentials in the Bethe-Salpeter equation. In practice this is difficult. One often utilizes a three-dimension reduction of the Bethe-Salpeter equation [84]. In the present work, following Ref. [73], we use the coupled-channel Kadyshevsky equation:

$$T_{\rho\rho'}^{\nu\nu',J}(\mathbf{p}', \mathbf{p}; \sqrt{s}) = V_{\rho\rho'}^{\nu\nu',J}(\mathbf{p}', \mathbf{p}) + \sum_{\rho'', \nu''} \int_0^\infty \frac{dp'' p''^2}{(2\pi)^3} \\ \times \frac{2\mu_{\nu''}^2 V_{\rho\rho''}^{\nu\nu'',J}(\mathbf{p}', \mathbf{p}'') T_{\rho''\rho'}^{\nu''\nu',J}(\mathbf{p}'', \mathbf{p}; \sqrt{s})}{(p''^2 + 4\mu_{\nu''}^2) \left(\sqrt{q_{\nu''}^2 + 4\mu_{\nu''}^2} - \sqrt{p''^2 + 4\mu_{\nu''}^2} + i\epsilon \right)}, \quad (18)$$

where \sqrt{s} is the total energy of the baryon-baryon system in the center-of-mass frame, $\mathbf{q}_{\nu''}$ is the relativistic on-shell momentum defined by $\sqrt{s} = \sqrt{M_{B_{1,\nu''}}^2 + \mathbf{q}_{\nu''}^2} + \sqrt{M_{B_{2,\nu''}}^2 + \mathbf{q}_{\nu''}^2}$, with $B_{1,\nu''}$ and $B_{2,\nu''}$ intermediate state baryons, and $\mu_{\nu''}$ is the reduced mass of the intermediate state. The labels ν, ν', ν'' denote the particle channels, and ρ, ρ', ρ'' denote the partial waves. Relativistic kinematics is used throughout to relate the laboratory momenta to the center-of-mass momenta.

To regularize the integration in the high-momentum region, baryon-baryon potentials are multiplied with an exponential form factor,

$$f_{\Lambda_F}(\mathbf{p}, \mathbf{p}') = \exp \left[-\left(\frac{\mathbf{p}}{\Lambda_F} \right)^{2n} - \left(\frac{\mathbf{p}'}{\Lambda_F} \right)^{2n} \right], \quad (19)$$

where $n = 2$ [85].

The Kadyshevsky equation is solved in particle basis in order to properly account for the physical thresholds and the Coulomb force in charged channels. The latter is treated with the Vincent-Phatak method [86].

3. Fitting procedure

In our approach, there are 12 LECs that need to be pinned down by fitting to the 36 YN scattering data as done in [73], which consist of 35 cross sections [87–90] and a $\Sigma^- p$ inelastic capture ratio at rest [91].

Due to the poor quality of experimental data, it is customary to consider the hypertriton ${}^3_\Lambda\text{H}$ binding energy [92, 93] as a further constraint, which is crucial in fixing the relative strength of the 1S_0 and 3S_1 contributions to Λp scattering. However, we are unable to perform a 3-body calculation at present, so we use as benchmarks the Λp S -wave scattering lengths extracted in the LO [56] and next-to-leading order (NLO) [58] calculations using the Weinberg approach, mainly because they combine to describe very well the hypertriton [94]. Empirically, the Λp 3S_1 contribution is more sensitive to the scattering data while the 1S_0 part is mainly decided by the hypertriton binding energy [95]. In addition, it seems necessary that $a_{1S_0}^{\Lambda p}$ should be neither smaller nor too larger than $a_{3S_1}^{\Lambda p}$, as shown in [96].

Another constraint that should be considered is the $\Sigma^+ p$ 3S_1 scattering length. A repulsive ΣN interaction with isospin $I = 3/2$ is obtained from recent experiments [97–103]. And the conventional G -matrix calculations [104] indicate that the 3S_1 partial wave for $I = 3/2$ ΣN should be at least moderately repulsive, therefore in our fits we require a positive $a_{3S_1}^{\Sigma^+ p}$.

Previous works in χEFT [56, 58, 73] showed that the optimum cutoff Λ_F may be around 600 MeV. Therefore we first tentatively fix Λ_F at 600 MeV. With this cutoff we find that the best description of the experimental data yield $a_{3S_1}^{\Lambda p} \approx -1.30 \pm 0.02$ fm and $a_{1S_0}^{\Lambda p} \approx -2.44^{+0.16}_{-0.54}$ fm. These numbers are between the LO and NLO Weinberg results, which are $a_{3S_1}^{\Lambda p} = -1.23$ fm (LO), $a_{3S_1}^{\Lambda p} = -1.54$ fm (NLO), $a_{1S_0}^{\Lambda p} = -1.91$ fm (LO), and $a_{1S_0}^{\Lambda p} = -2.91$ fm (NLO). Best fits within $\Lambda_F = 500 - 850$ MeV yield similar scattering lengths. In the results presented below, we fix $a_{3S_1}^{\Lambda p} = -1.32$ fm and $a_{1S_0}^{\Lambda p} = -2.44$ fm.¹ It should be noted that at present we could in principle choose other combinations within the uncertainties allowed in the best fits. To fix them uniquely, more experimental inputs are needed.

We have made an attempt at the combined fit to the NN and YN data, in which strict SU(3) symmetry was imposed upon the contact terms so that no additional LECs are needed. However, we failed to describe the NN and YN data simultaneously. As a result, consistent with previous NLO results in the Weinberg approach [58], we conclude that one needs to treat more carefully the SU(3) symmetry breaking in order to simultaneously describe both the NN and the YN systems in χEFT .

4. Results and discussions

With the three additional constraints $a_{1S_0}^{\Lambda p} = -2.44$ fm, $a_{3S_1}^{\Lambda p} = -1.32$ fm and $a_{3S_1}^{\Sigma^+ p} > 0$ as explained above, we perform a fit to

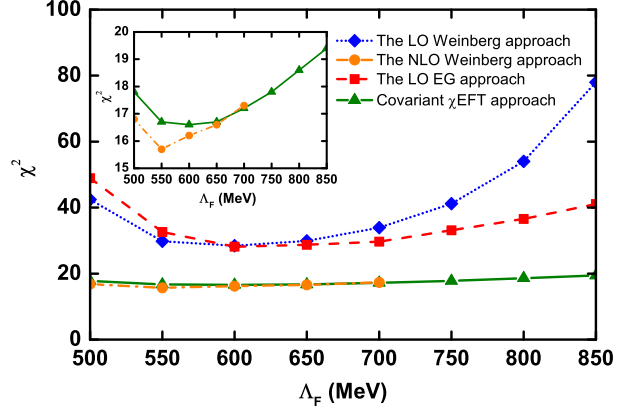


Figure 2: χ^2 as a function of the cutoff in the Weinberg approach at LO (blue dotted line), NLO (orange dashed-dotted line) [58], the EG approach at LO (red dashed line) and the covariant χEFT approach at LO (green solid line).

the 36 scattering data while varying the cutoff Λ_F . The dependence of the χ^2 on the cutoff is shown in Fig. 2, in comparison with other approaches. One can see that our new covariant χEFT approach shows a clear improvement in describing the YN data compared with the Weinberg and EG approach at LO, and the cutoff dependence is much mitigated, both of which are comparable with the NLO Weinberg approach [58], which, however, has 23 LECs. The minimum value of the χ^2 is about 16.7, located at $\Lambda_F = 550 - 650$ MeV. Note that the NSC97a-f [29] models, which provide the best description of the 36 scattering data among the phenomenological potentials, also have a χ^2 around 16.

The best fitted LECs obtained with $\Lambda_F = 600$ MeV are listed in Table 1. Since the LECs in the Λp 1S_0 partial wave can not be uniquely determined as mentioned previously, we only show a typical case here. One should note that these LECs are certain combinations of those appearing in the Lagrangians, and hence they are not necessarily of the same order of magnitude (see, e.g., Refs. [56, 58]).

In Fig. 3 we compare the descriptions of the experimental cross sections that we have used in the fitting procedure with the LO Weinberg approach. The NSC97f [29] and Jülich 04 results [34] are also shown for comparison. It is clear that the covariant χEFT approach can better reproduce the experimental data points, in particular for $\Lambda p \rightarrow \Lambda p$ and $\Sigma^+ p \rightarrow \Sigma^+ p$.

Due to the lack of near threshold experimental data, the value of $\Lambda p \rightarrow \Lambda p$ cross section at rest is not yet known. Our value is about 350 mb, which is smaller than the two phenomenological models. Our result in the $\Sigma^- p \rightarrow \Lambda n$ reaction is similar to the LO Weinberg approach and NSC97f results, but quite different from the Jülich 04 model. This channel can partially reflect the nature of $\Lambda N - \Sigma N$ coupling, which is crucial in hypernuclear structure calculations [18]. It is interesting to note that the Jülich 04 model predicts an overbound Λ single particle potential $U_\Lambda(0)$ in G -matrix calculations. On the other hand, the results from the former two are much closer to the empirical value, c.f. [104] and references therein.

One should note that the improved description of the scatter-

¹We have chosen a larger $a_{3S_1}^{\Lambda p}$ given the fact that most phenomenological studies seem to prefer a larger scattering length in this channel.

Table 1: Low-energy constants (in units of 10^4 GeV^{-2}) at $\Lambda_F = 600 \text{ MeV}$ in the covariant χEFT approach.

LECs	$C_{1S0}^{\Lambda\Lambda}$	$C_{1S0}^{\Sigma\Sigma}$	$C_{3S1}^{\Lambda\Lambda}$	$C_{3S1}^{\Sigma\Sigma}$	$C_{3S1}^{\Lambda\Sigma}$	$\hat{C}_{1S0}^{\Lambda\Lambda}$	$\hat{C}_{1S0}^{\Sigma\Sigma}$	$\hat{C}_{3S1}^{\Lambda\Lambda}$	$\hat{C}_{3S1}^{\Sigma\Sigma}$	$\hat{C}_{3S1}^{\Lambda\Sigma}$	$C_{3P1}^{\Lambda\Lambda}$	$C_{3P1}^{\Sigma\Sigma}$
	-0.0223	-0.0528	0.0032	0.0842	0.0232	3.9185	4.0681	0.4190	-0.4132	0.7326	0.2044	0.2616

ing data by the covariant χEFT scheme for the most part arises from the contact terms. In the LO Weinberg approach, contact terms only appear in central and spin-spin potentials without momentum dependence, which only contribute to the 1S_0 and 3S_1 partial waves. In the covariant power counting, tensor, spin-orbit, asymmetric spin-orbit and quadratic spin-orbit terms appear at LO in addition to the central and spin-spin terms of the Weinberg approach, namely the momentum dependent terms of $R_p^2(p')$ in Eqs. (7-14). These terms are responsible for the improved description. Meanwhile, relativistic corrections in the one-pseudoscalar-meson-exchange potentials contribute very little. These are similar to the NN sector. A more detailed discussion can be found in Ref. [83].

5. Summary and outlook

We have studied strangeness $S = -1$ hyperon-nucleon scattering at leading order in a covariant framework of chiral effective field theory. Starting from the covariant chiral Lagrangian, the small components of the baryon spinors are retained in deriving the potentials in order to preserve Lorentz invariance. Strict $SU(3)$ symmetry is imposed upon contact terms, which yield 12 independent low energy constants. $SU(3)$ symmetry is broken in the one-pseudoscalar-meson-exchange potentials because of the mass difference of exchanged mesons. The potentials are iterated using the Kadyshevsky equation. A quite satisfactory description of the 36 hyperon-nucleon scattering data is obtained and the cutoff dependence is shown to be mitigated, both of which are comparable with the next-to-leading order Weinberg approach. However, the latter has 11 more low energy constants. The relativistic interaction obtained is shown to be more efficient than the Weinberg approach and may provide essential inputs to relativistic hypernuclear structure studies.

Acknowledgements

This work is partly supported by the National Natural Science Foundation of China under Grants No. 11375024, No. 11522539, and No. 11375120, the China Postdoctoral Science Foundation under Grant No. 2016M600845, and the Fundamental Research Funds for the Central Universities.

References

References

- [1] M. Gell-Mann, Phys. Rev. **92**, 833 (1953).
- [2] T. Nakano and K. Nishijima, Prog. Theor. Phys. **10**, 581 (1953).
- [3] M. Danysz, and J. Pniewski, Philos. Mag. Ser. 5 **44**, 348 (1953).
- [4] M. Bedjidian *et al.* [CERN-Lyon-Warsaw Collaboration], Phys. Lett. **83B**, 252 (1979).

- [5] R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977) Erratum: [Phys. Rev. Lett. **38**, 617 (1977)].
- [6] A. Nogga, H. Kamada and W. Glöckle, Phys. Rev. Lett. **88**, 172501 (2002) [nucl-th/0112060].
- [7] A. Esser *et al.* [A1 Collaboration], Phys. Rev. Lett. **114**, 232501 (2015) [arXiv:1501.06823 [nucl-ex]].
- [8] T. O. Yamamoto *et al.* [J-PARC E13 Collaboration], Phys. Rev. Lett. **115**, 222501 (2015) [arXiv:1508.00376 [nucl-ex]].
- [9] D. Gazda and A. Gal, Phys. Rev. Lett. **116**, 122501 (2016) [arXiv:1512.01049 [nucl-th]].
- [10] S. R. Beane *et al.* [NPLQCD Collaboration], Phys. Rev. Lett. **106**, 162001 (2011) [arXiv:1012.3812 [hep-lat]].
- [11] T. Inoue *et al.* [HAL QCD Collaboration], Phys. Rev. Lett. **106**, 162002 (2011) [arXiv:1012.5928 [hep-lat]].
- [12] J. Haidenbauer and U. -G. Meißner, Phys. Lett. B **706**, 100 (2011) [arXiv:1109.3590 [hep-ph]].
- [13] J. K. Ahn *et al.* [E373 (KEK-PS) Collaboration], Phys. Rev. C **88**, 014003 (2013).
- [14] Y. Yamaguchi and T. Hyodo, arXiv:1607.04053 [hep-ph].
- [15] K. Nakazawa *et al.*, PTEP **2015**, no. 3, 033D02 (2015).
- [16] F. Hiruma, http://j-parc.jp/researcher/Hadron/en/pac_1101/pdf/KEK-J-PARC-PAC20110101.pdf.
- [17] F. Hauenstein *et al.* [COSY-TOF Collaboration], arXiv:1607.04783 [nucl-ex].
- [18] E. Hiyama, S. Ohnishi, B. F. Gibson and T. A. Rijken, Phys. Rev. C **89**, 061302 (2014) [arXiv:1405.2365 [nucl-th]].
- [19] X. R. Zhou, H.-J. Schulze, H. Sagawa, C. X. Wu and E. G. Zhao, Phys. Rev. C **76**, 034312 (2007).
- [20] E. Massot, J. Margueron and G. Chanfray, Europhys. Lett. **97**, 39002 (2012) [arXiv:1201.2772 [nucl-th]].
- [21] H.-J. Schulze and T. Rijken, Phys. Rev. C **84**, 035801 (2011).
- [22] J. N. Hu, A. Li, H. Toki and W. Zuo, Phys. Rev. C **89**, 025802 (2014) [arXiv:1307.4154 [nucl-th]].
- [23] T. Miyatsu, S. Yamamuro and K. Nakazato, Astrophys. J. **777**, 4 (2013) [arXiv:1308.6121 [astro-ph.HE]].
- [24] R. Mallick, Phys. Rev. C **87**, 025804 (2013) [arXiv:1207.4872 [astro-ph.HE]].
- [25] P. Demorest, T. Pennucci, S. Ransom, M. Roberts and J. Hessels, Nature **467**, 1081 (2010) [arXiv:1010.5788 [astro-ph.HE]].
- [26] J. Antoniadis *et al.*, Science **340**, 6131 (2013) [arXiv:1304.6875 [astro-ph.HE]].
- [27] M. M. Nagels, T. A. Rijken and J. J. de Swart, Phys. Rev. D **15**, 2547 (1977).
- [28] P. M. M. Maessen, T. A. Rijken and J. J. de Swart, Phys. Rev. C **40**, 2226 (1989).
- [29] T. A. Rijken, V. G. J. Stoks and Y. Yamamoto, Phys. Rev. C **59**, 21 (1999) [nucl-th/9807082].
- [30] T. A. Rijken and Y. Yamamoto, Phys. Rev. C **73**, 044008 (2006) [nucl-th/0603042].
- [31] M. M. Nagels, T. A. Rijken and Y. Yamamoto, arXiv:1501.06636 [nucl-th].
- [32] B. Holzenkamp, K. Holinde and J. Speth, Nucl. Phys. A **500**, 485 (1989).
- [33] A. Reuber, K. Holinde and J. Speth, Nucl. Phys. A **570**, 543 (1994).
- [34] J. Haidenbauer and U. -G. Meißner, Phys. Rev. C **72**, 044005 (2005) [nucl-th/0506019].
- [35] U. Straub, Z. Y. Zhang, K. Brauer, A. Faessler, S. B. Khadkikar and G. Lubeck, Nucl. Phys. A **483**, 686 (1988).
- [36] U. Straub, Z. Y. Zhang, K. Brauer, A. Faessler, S. B. Khadkikar and G. Lubeck, Nucl. Phys. A **508**, 385C (1990).
- [37] Z. Y. Zhang, A. Faessler, U. Straub and L. Y. Glozman, Nucl. Phys. A **578**, 573 (1994).
- [38] Z. Y. Zhang, Y. W. Yu, P. N. Shen, L. R. Dai, A. Faessler and U. Straub, Nucl. Phys. A **625**, 59 (1997).
- [39] J. L. Ping, F. Wang and J. T. Goldman, Nucl. Phys. A **657**, 95 (1999)

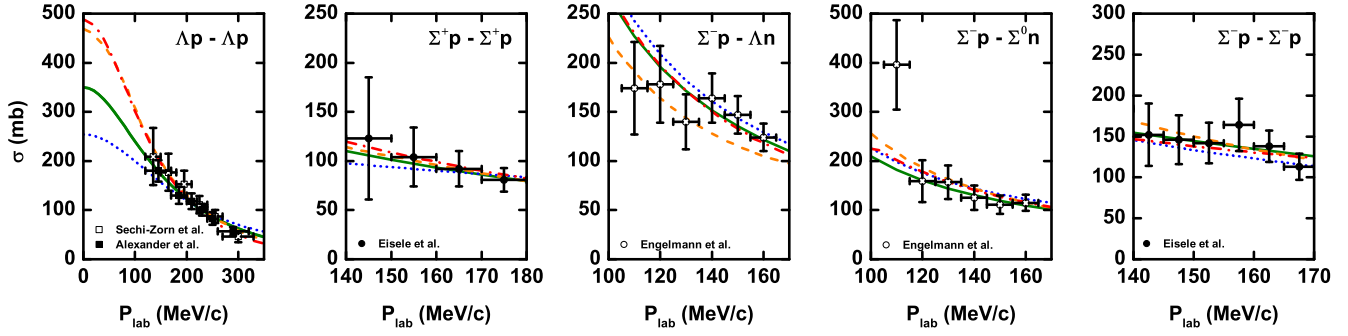


Figure 3: Cross sections in the Weinberg approach (blue dotted lines) and covariant χ EFT approach (green solid lines) at LO as functions of the laboratory momentum at $\Lambda_F = 600$ MeV. For reference, the NSC97f [29] (red dash-dotted lines) and Jülich 04 [34] (orange dashed lines) results are also shown. The experimental data are taken from Sechi-Zorn *et al.* [87], Alexander *et al.* [88], Eisele *et al.* [90], and Engelmann *et al.* [89].

- [nucl-th/9812068].
- [40] Y. Fujiwara, C. Nakamoto and Y. Suzuki, Phys. Rev. Lett. **76**, 2242 (1996).
- [41] Y. Fujiwara, Y. Suzuki and C. Nakamoto, Prog. Part. Nucl. Phys. **58**, 439 (2007) [nucl-th/0607013].
- [42] S. R. Beane *et al.* [NPLQCD Collaboration], Nucl. Phys. A **794**, 62 (2007) [hep-lat/0612026].
- [43] H. Nemura, N. Ishii, S. Aoki and T. Hatsuda, Phys. Lett. B **673**, 136 (2009) [arXiv:0806.1094 [nucl-th]].
- [44] S. R. Beane *et al.* [NPLQCD Collaboration], Phys. Rev. D **81**, 054505 (2010) [arXiv:0912.4243 [hep-lat]].
- [45] T. Inoue *et al.* [HAL QCD Collaboration], Prog. Theor. Phys. **124**, 591 (2010) [arXiv:1007.3559 [hep-lat]].
- [46] S. R. Beane *et al.* [NPLQCD Collaboration], Phys. Rev. D **85**, 054511 (2012) [arXiv:1109.2889 [hep-lat]].
- [47] K. Sasaki *et al.* [HAL QCD Collaboration], PTEP **2015**, 113B01 (2015) [arXiv:1504.01717 [hep-lat]].
- [48] T. Doi *et al.*, arXiv:1512.01610 [hep-lat].
- [49] T. Doi *et al.*, arXiv:1512.04199 [hep-lat].
- [50] P. F. Bedaque and U. van Kolck, Ann. Rev. Nucl. Part. Sci. **52**, 339 (2002) [nucl-th/0203055].
- [51] E. Epelbaum, H. W. Hammer and U. -G. Meißner, Rev. Mod. Phys. **81**, 1773 (2009) [arXiv:0811.1338 [nucl-th]].
- [52] R. Machleidt and D. R. Entem, Phys. Rept. **503**, 1 (2011) [arXiv:1105.2919 [nucl-th]].
- [53] S. Weinberg, Phys. Lett. B **251**, 288 (1990).
- [54] S. Weinberg, Nucl. Phys. B **363**, 3 (1991).
- [55] X. W. Kang, J. Haidenbauer and U. -G. Meißner, JHEP **1402**, 113 (2014) [arXiv:1311.1658 [hep-ph]].
- [56] H. Polinder, J. Haidenbauer and U. -G. Meißner, Nucl. Phys. A **779**, 244 (2006) [nucl-th/0605050].
- [57] J. Haidenbauer, U. -G. Meißner, A. Nogga and H. Polinder, Lect. Notes Phys. **724**, 113 (2007) [nucl-th/0702015 [NUCL-TH]].
- [58] J. Haidenbauer, S. Petschauer, N. Kaiser, U. -G. Meißner, A. Nogga and W. Weise, Nucl. Phys. A **915**, 24 (2013) [arXiv:1304.5339 [nucl-th]].
- [59] H. Polinder, J. Haidenbauer and U. -G. Meißner, Phys. Lett. B **653**, 29 (2007) [arXiv:0705.3753 [nucl-th]].
- [60] J. Haidenbauer and U. -G. Meißner, Phys. Lett. B **684**, 275 (2010) [arXiv:0907.1395 [nucl-th]].
- [61] J. Haidenbauer, U. -G. Meißner and S. Petschauer, Nucl. Phys. A **954**, 273 (2016) [arXiv:1511.05859 [nucl-th]].
- [62] G. P. Lepage, nucl-th/9706029.
- [63] M. C. Birse, Phys. Rev. C **74**, 014003 (2006) [nucl-th/0507077].
- [64] A. Nogga, R. G. E. Timmermans and U. van Kolck, Phys. Rev. C **72**, 054006 (2005) [nucl-th/0506005].
- [65] E. Epelbaum and U. -G. Meißner, Few Body Syst. **54**, 2175 (2013) [nucl-th/0609037].
- [66] B. Long and U. van Kolck, Annals Phys. **323**, 1304 (2008) [arXiv:0707.4325 [quant-ph]].
- [67] C. -J. Yang, C. Elster and D. R. Phillips, Phys. Rev. C **80**, 044002 (2009) [arXiv:0905.4943 [nucl-th]].
- [68] M. P. Valderrama, Phys. Rev. C **83**, 024003 (2011) [arXiv:0912.0699 [nucl-th]].
- [69] B. Long and C. J. Yang, Phys. Rev. C **85**, 034002 (2012) [arXiv:1111.3993 [nucl-th]].
- [70] E. Epelbaum and J. Gegelia, Phys. Lett. B **716**, 338 (2012) [arXiv:1207.2420 [nucl-th]].
- [71] E. Epelbaum, A. M. Gasparyan, J. Gegelia and H. Krebs, Eur. Phys. J. A **51**, 71 (2015) [arXiv:1501.01191 [nucl-th]].
- [72] V. G. J. Stoks, R. A. M. Klomp, M. C. M. Rentmeester and J. J. de Swart, Phys. Rev. C **48**, 792 (1993).
- [73] K. -W. Li, X. -L. Ren, L. S. Geng and B. Long, Phys. Rev. D **94**, 014029 (2016) [arXiv:1603.07802 [hep-ph]].
- [74] L. S. Geng, J. Martin Camalich, L. Alvarez-Ruso and M. J. Vicente Vacas, Phys. Rev. Lett. **101**, 222002 (2008) [arXiv:0805.1419 [hep-ph]].
- [75] L. S. Geng, J. Martin Camalich and M. J. Vicente Vacas, Phys. Rev. D **79**, 094022 (2009) [arXiv:0903.4869 [hep-ph]].
- [76] L. S. Geng, X. -L. Ren, J. Martin-Camalich and W. Weise, Phys. Rev. D **84**, 074024 (2011) [arXiv:1108.2231 [hep-ph]].
- [77] X. -L. Ren, L. S. Geng, J. Martin Camalich, J. Meng and H. Toki, JHEP **1212**, 073 (2012) [arXiv:1209.3641 [nucl-th]].
- [78] X. -L. Ren, L. -S. Geng and J. Meng, Phys. Rev. D **91**, 051502 (2015) [arXiv:1404.4799 [hep-ph]].
- [79] L. S. Geng, N. Kaiser, J. Martin-Camalich and W. Weise, Phys. Rev. D **82**, 054022 (2010) [arXiv:1008.0383 [hep-ph]].
- [80] L. S. Geng, M. Altenbuchinger and W. Weise, Phys. Lett. B **696**, 390 (2011) [arXiv:1012.0666 [hep-ph]].
- [81] M. Altenbuchinger, L. S. Geng and W. Weise, Phys. Lett. B **713**, 453 (2012) [arXiv:1109.0460 [hep-ph]].
- [82] L. S. Geng, Front. Phys. (Beijing) **8**, 328 (2013) [arXiv:1301.6815 [nucl-th]].
- [83] X. -L. Ren, K. -W. Li, L. S. Geng, B. -W. Long, P. Ring and J. Meng, arXiv:1611.08475 [nucl-th].
- [84] R. M. Woloshyn and A. D. Jackson, Nucl. Phys. B **64**, 269 (1973).
- [85] E. Epelbaum, W. Glöckle and U. -G. Meißner, Nucl. Phys. A **747**, 362 (2005) [nucl-th/0405048].
- [86] C. M. Vincent and S. C. Phatak, Phys. Rev. C **10**, 391 (1974).
- [87] B. Sechi-Zorn, B. Kehoe, J. Twitty and R. A. Burnstein, Phys. Rev. **175**, 1735 (1968).
- [88] G. Alexander, U. Karshon, A. Shapira, G. Yekutieli, R. Engelmann, H. Filthuth and W. Lughofer, Phys. Rev. **173**, 1452 (1968).
- [89] R. Engelmann, H. Filthuth, V. Hepp, E. Kluge, Phys. Lett. **21**, 587 (1966).
- [90] F. Eisele, H. Filthuth, W. Foehlich, V. Hepp and G. Zech, Phys. Lett. B **37**, 204 (1971).
- [91] V. Hepp and H. Schleich, Z. Phys. **214**, 71 (1968).
- [92] M. Juric *et al.*, Nucl. Phys. B **52**, 1 (1973).
- [93] D. H. Davis, AIP Conf. Proc. **224**, 38 (1991).
- [94] A. Nogga, Nucl. Phys. A **914**, 140 (2013).
- [95] C. L. Korpa, A. E. L. Dieperink and R. G. E. Timmermans, Phys. Rev. C **65**, 015208 (2002) [nucl-th/0109072].
- [96] K. Tominaga, T. Ueda, M. Yamaguchi, N. Kijima, D. Okamoto,

- K. Miyagawa and T. Yamada, Nucl. Phys. A **642**, 483 (1998).
- [97] C. J. Batty, E. Friedman and A. Gal, Phys. Lett. B **335**, 273 (1994).
- [98] J. Mares, E. Friedman, A. Gal and B. K. Jennings, Nucl. Phys. A **594**, 311 (1995) [nucl-th/9505003].
- [99] S. Bart *et al.*, Phys. Rev. Lett. **83**, 5238 (1999).
- [100] H. Noumi *et al.*, Phys. Rev. Lett. **89**, 072301 (2002) Erratum: [Phys. Rev. Lett. **90**, 049902 (2003)].
- [101] P. K. Saha *et al.*, Phys. Rev. C **70**, 044613 (2004) [nucl-ex/0405031].
- [102] M. Kohno, Y. Fujiwara, Y. Watanabe, K. Ogata and M. Kawai, Phys. Rev. C **74**, 064613 (2006) [nucl-th/0611080].
- [103] J. Dabrowski and J. Rozynek, Phys. Rev. C **78**, 037601 (2008).
- [104] J. Haidenbauer and U.-G. Meißner, Nucl. Phys. A **936**, 29 (2015) [arXiv:1411.3114 [nucl-th]].